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Theoretical Aspects of Double Beta Decay

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ABSTRACT

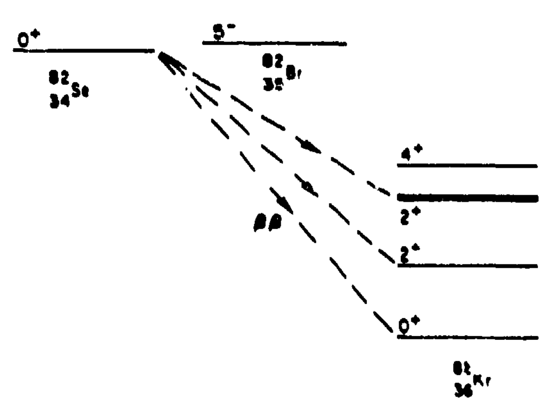
A summary of recent work in the theory of double beta decay is given.

1. Introduction

Considerable effort has been expended recently in theoretical studies of double beta decay. Much of this work has focussed on the constraints this process places on gauge theories of the weak interaction, in general, and on the neutrino mass matrix, in particular. In addition, interesting nuclear structure questions have arisen in studies of double beta decay matrix elements. After briefly reviewing the theory of double beta decay, I will summarize some of the progress that has been made in these areas.

Double beta decay is the process by which a parent nucleus (A,Z) can decay to a daughter (A,Z±2) by emitting two electrons or positrons. This process can be observed because the nuclear pairing force increases the binding energy of nuclei with even numbers of protons and neutrons relative to those with odd numbers. Consequently the ordinary

Figure 1: ββ decay scheme of ^{82}Se



β decay of an even-even nucleus $(A,Z)\to(A,Z\pm1)$ is frequently forbidden energetically, leaving $\beta\beta$ decay as the only allowed decay mode.

The question historically associated with $\beta\beta$ decay is whether the electron neutrino should be described by a Dirac field ($\nu_e \perp \bar{\nu}_e$) or by a Majorana field ($\nu_e=\bar{\nu}_e$). As any fermion having a charge or a magnetic moment necessarily has a distinct antiparticle, the neutrino is unique in permitting these alternative descriptions. Prior to 1957 it was believed that $\beta\beta$ decay experiments had determined which of these descriptions was correct. We define the electron neutrino and its antiparticle by

$$n\to p+e^-+\bar{\nu}_e \qquad (1a)$$

$$\nu_e+n\to p+e^- \qquad (1b)$$

It follows that the decay $(A,Z)\to(A,Z+2)$ can occur by successive β decays (Eq. (1a)) with virtual excitation of the intermediate nucleus $(A,Z+1)$, as shown in Fig. 2a:

$$2n\to n+p+e^-+\bar{\nu}_e\to 2p+2e^-+2\bar{\nu}_e \qquad (2)$$

The occurrence of this 2ν sequence does not depend on the charge conjugation properties of the neutrino. A second decay mode (first discussed by Racah[2]) will occur only if the neutrino is a Majorana particle

$$2n\to n+p+e^-+\bar{\nu}_e$$
$$\equiv n+p+e^-+\nu_e\to 2p+2e^- \qquad (3)$$

producing a neutrinoless final state, as shown in Fig. 2b.

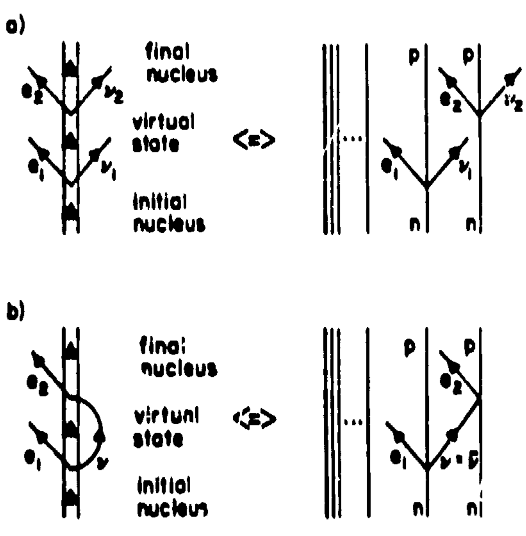


Figure 2: Two-nucleon mechanisms for (a) two-neutrino and (b) no-neutrino $\beta\beta$ decay

The prejudice for a conserved lepton number grew out of the observation that, for a Majorana neutrino, the expected rates for the processes shown in Fig. (2) are

$$\omega^{2\nu} \sim (10^{-20}\text{-}10^{-24})/y$$
$$\omega^{0\nu} \sim (10^{-13}\text{-}10^{-15})/y \ ,$$

reflecting the phase space difference between 2ν and 0ν $\beta\beta$ decay. By 1951 a series of geochemical, counter, and radiochemical experiments had established

$$\omega^{exp} \lesssim 10^{-18}/y$$

leading to the conclusion that the neutrino must be a Dirac particle to explain the absence of 0ν decay. Lepton number was introduced as the quantum number distinguishing the neutrino from its antiparticle, $\ell(\nu_e)=+1$ and $\ell(\bar{\nu}_e)=-1$.

The experimental verification of apparent maximal parity violation[3] in β decay exposed a flaw in this argument. The neutrino and antineutrino participating in Eqs. (1) are left-handed and right-handed, respectively:

$$n \to p + e^- + \bar{\nu}_e^{RH} \qquad (1a')$$
$$\nu_e^{LH} + n \to p + e^- \qquad (1b')$$

Thus the Racah sequence for a Majorana neutrino

$$2n \to n + p + e^- + \bar{\nu}^{RH}$$
$$\equiv n + p + e^- + \nu_e^{RH} \nrightarrow 2p + 2e^-$$

is forbidden because the right-handed neutrino has the wrong helicity to complete the last step. Therefore the absence of this decay mode implies neither a Dirac neutrino nor a conserved lepton number.

The intense modern interest in double beta decay stems from the expectation in grand unified theories that this "γ_5-invariance" of the weak leptonic current is, in fact, only approximate. Because of the favorable phase space for 0ν $\beta\beta$ decay, experimental bounds on this process can place very stringent constraints on the masses and right-handed couplings of possible Majorana neutrinos. For instance, some GUTS predict Majorana masses of the form[4]

$$m_\nu \sim \frac{M_D^2}{M_R}$$

where the Dirac mass M_D may be that of a quark or charged lepton. The experimental talks in this session will describe efforts to probe masses $m_\nu \sim 1$ eV in modern searches for 0ν $\beta\beta$ decay. Thus, for $M_D \sim$ (1 MeV - 1 GeV), these experiments will test mass scales $M_R \sim (10^3\text{-}10^9)$ GeV, a range that begins at the highest energies directly accessible in modern accelerators.

2. Two-neutrino $\beta\beta$ Decay in the Standard Model

I begin with a brief description of the 2ν $\beta\beta$ decay process shown in Fig. 2a. Two approximations are frequently made by theorists:

i) Each nuclear β decay is evaluated in the allowed approximation where only the Fermi and Gamow-Teller operators ($\tau_+(i)$ and $\vec{\sigma}(i)\tau_+(i)$) are retained;

ii) The sum over virtual intermediate nuclear states is performed by closure after replacing the nuclear excitation energy appearing in the energy denominator by an average value.

The first approximation restricts the states populated in the daughter nucleus by the decay of a 0^+ parent to those with $J^\pi=0^+,1^+$, and 2^+. In fact, decays to 1^+ and 2^+ states are strongly suppressed. Decays between 0^+ states are mediated by two matrix elements

$$M_{GT} = \langle 0_f^+ | \tfrac{1}{2} \sum_{ij} \vec{\sigma}(i)\cdot\vec{\sigma}(j)\tau_+(i)\tau_+(j) | 0_i^+ \rangle \tag{4a}$$

$$M_F = \langle 0_f^+ | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j) | 0_i^+ \rangle \tag{4b}$$

The double Fermi matrix element vanishes in the limit of good isospin and quite generally is small. Thus the decay rate can be written approximately as

$$\omega_{2\nu} \sim f_{GT} \, |M_{GT}|^2 \tag{5}$$

where the phase space factor f_{GT} is

$$f_{GT} \sim \frac{2m_e^{11}(G\cos\theta_c)^4}{\pi^7 7!} \quad \frac{F^2(z)}{\langle E\rangle^2} \quad [\tilde{T}_o^7 + \dots + \frac{\tilde{T}_o^{11}}{1980}]$$

with θ_c the Cabibbo angle, $F(z)$ a correction for distortions of the electron plane waves in the Coulomb field of the nucleus, <E> the average intermediate state excitation energy, and $\tilde{T}_o$ the total kinetic energy carried off by the leptons in units of $m_e c^2$.

3. Neutrinoless ββ decay

I use the following effective β decay Lagrangian density

$$L(x) = \frac{G\cos\theta_c}{\sqrt{2}} [\bar{\psi}_e(x)\gamma^\mu(1-\gamma_5)\psi_{\nu_L}(x)\bar{u}(x)\gamma_\mu((1-\gamma_5)+\eta_{LR}(1+\gamma_5))d(x)$$
$$+ \bar{\psi}_e(x)\gamma^\mu(1+\gamma_5)\psi'_{\nu_R}(x)\bar{u}(x)\gamma_\mu(\eta_{RR}(1+\gamma_5)+\eta_{RL}(1-\gamma_5))d(x)]$$
$$+ L_m(x) \tag{6}$$

with $\psi_e(x)$, $\psi_{\nu_L}(x)$, $\psi'_{\nu_R}(x)$, $u(x)$, and $d(x)$, the electron, left- and right-handed neutrino, and the u and d quark fields. The usual V-A Lagrangian of the standard model is obtained by setting $\eta_{LR}=\eta_{RR}=\eta_{RL}=0$. Note that η_{LR}, η_{RR}, and η_{RL} are defined so that the first subscript refers to the leptonic current (left- or right-handed) and the second to the hadronic current. The right-handed couplings η_{RR} and η_{RL} explicitly break the γ_5-invariance of the weak leptonic current.

The left-and right-handed current neutrino fields can be written in terms of 2n Majorana mass eigenstate fields $v_i(x)$, where n is the number of generations (e,μ,τ, ...) in some underlying theory:

$$\psi_{\nu_L} = \sum_{i=1}^{2n} U^L_{ei} v_i \qquad \psi'_{\nu_R} = \sum_{i=1}^{2n} U^R_{ei} v_i \tag{7}$$

The coefficients U^L_{ei} and U^R_{ei} that express the current fields in terms of the mass eigenstate fields can be determined by diagonalizing the 4n×4n mass term

$$L_m(x) = -\tfrac{1}{2}(\overline{\psi^c_L}, \overline{\psi'_R}, \overline{\psi_L}, \overline{\psi'^c_R}) \begin{pmatrix} 0 & 0 & M_L & M_D^T \\ 0 & 0 & M_D & M_R^\dagger \\ M_L^\dagger & M_D^\dagger & 0 & 0 \\ M_D^* & M_R & 0 & 0 \end{pmatrix} \begin{pmatrix} \psi^c_L \\ \psi'_R \\ \psi_L \\ \psi'^c_R \end{pmatrix} \tag{8}$$

where $\psi^c = C\psi C^{-1}$ with C the charge conjugation operator. The Majorana fields $v_i(x)$ can be written in terms of four-component Dirac spinors as

$$v_i(\vec{x},0) = \int \frac{d^3p}{(2\pi)^{3/2}} \sqrt{\frac{m_\nu}{p_o}} \sum_{\pm s} [e_i(p,s)U(p,s)e^{i\vec{p}\cdot\vec{x}}$$

$$+ \lambda_i^c e_i^\dagger(p,s)V(p,s)e^{-i\vec{p}\cdot\vec{x}}] \tag{9}$$

The following approximations are made in deriving the decay rate:

(1) Only the Fermi and Gamow-Teller nuclear β decay operators and electron s- and p-waves are retained.

(2) The closure approximation is invoked in the sum over virtual nuclear states. This approximation is better justified than in 2ν decay because of the presence of an energetic virtual neutrino in the intermediate state.

(3) L(x) is taken to be CP-invariant.

(4) All neutrino mass eigenstates are assumed to be light ($\lesssim$ 10 MeV) on the scale of nuclear energies.

The last two approximations are made to simplify this presentation.

As a consequence of (4) one finds that the 0ν rate formulae depend on separate particle physics and nuclear physics quantities. The rate depends quadratically on the lepton-number-violating masses and couplings

$$\langle m_\nu\rangle_{LL} = \sum_{i=1}^{2n} \lambda_i^{CP} |U_{ei}^L|^2 m_i \tag{10a}$$

$$\eta_{RL}\langle 1\rangle_{LR} \text{ with } \langle 1\rangle_{LR} = \sum_{i=1}^{2n} \lambda_i^{CP} U_{ei}^L U_{ei}^R \tag{10b}$$

$$\eta_{RR} \langle 1\rangle_{LR} \tag{10c}$$

$$\eta_{RR}^2 \langle m_\nu\rangle_{RR} \equiv \eta_{RR}^2 \sum_{i=1}^{2n} \lambda_i^{CP} |U_{ei}^R|^2 m_i \tag{10d}$$

with $m_i \geq 0$ and with $\lambda_i^{CP} = \pm\, i$ the CP eigenvalue of the mass eigenstate

$(CPe_i(p,s)(CP)^{-1} = \lambda_i^{CP} e(-p,s))$. Note that the mass terms involve either left-handed or right-handed neutrino fields, while the right-handed coupling terms depend on the interferences $U^L_{ei} U^R_{ei}$ between left- and right-handed neutrinos.

Experimental limits on 0ν $\beta\beta$ decay constrain the strength of the right-handed current couplings and the form of the neutrino mass matrix. Several limiting forms of the mass matrix have interesting consequences for $\beta\beta$ decay:

(1) The Dirac limit, $M_L = M_R = 0$, $M_D \neq 0$. The mass matrix diagonalization yields pairs of degenerate Majorana mass eigenstates with opposite CP. These pairs cancel explicitly in the expressions for $\langle m_\nu \rangle_{LL}$, $\langle m_\nu \rangle_{RR}$, and $\langle 1 \rangle_{LR}$. As it must, the rate for 0ν $\beta\beta$ decay vanishes in the Dirac limit.

(2) The pseudo-Dirac limit,[5] which I define as $M_L = M_R \ll M_D$ with, for simplicity, n=1. Two nearly degenerate Majorana mass eigenstates $m_i = |M_D \pm M_L|$ with opposite CP result. One finds $\langle m_\nu \rangle^2_{LL} = \langle m_\nu \rangle^2_{RR} = M_L^2$. Thus the relevant mass in $\beta\beta$ decay is the small mass splitting of the mass eigenstates.

(3) The Gell-Mann, Ramond, Slansky[4] limit, $M_L = 0$, $M_D \neq 0$, $|M_R| \gg |M_D|$, with n=1 for simplicity. One light and one heavy mass eigenstate result. These contribute dominantly to ψ_{ν_L} and ψ'_{ν_R}, respectively The light mass contribution to $\langle m_\nu \rangle_{LL}$ is proportional to M_D^2/M_R.

The 0ν $\beta\beta$ decay rate also depends quadratically on a variety of nuclear matrix elements. For $0^+ \to 0^+$ decay these are

$$\langle 0_f^+ | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\, \vec{\sigma}(i)\cdot\vec{\sigma}(j)\, \frac{g(r_{ij})}{r_{ij}} \,|0^+\rangle \tag{11a}$$

$$\langle 0^+ | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\, \frac{g(r_{ij})}{r_{ij}} \,|0^+\rangle \tag{11b}$$

$$\langle 0_f^+ | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\hat{r}_{ij}\cdot\vec{\sigma}(i)\hat{r}_{ij}\cdot\vec{\sigma}(j)\, \frac{g(r_{ij})}{r_{ij}} \,|0_i^+\rangle \tag{11c}$$

$$\langle 0_f^+ | \tfrac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\hat{R}_{ij}\cdot(\hat{r}_{ij}\times(\vec{\sigma}(i)-\vec{\sigma}(j)))\, \frac{g(r_{ij})R_{ij}}{r_{ij}^2} \,|0^+\rangle \tag{11d}$$

where $\vec{r}_{ij} = \vec{r}_i - \vec{r}_j$ and $\vec{R}_{ij} = \vec{r}_i + \vec{r}_j$. One can take $g(r_{ij}) \sim 1$ for light neutrinos ($m_\nu << 1/R_o$, with R_o the nuclear radius). As only the first two matrix elements contribute to terms in the decay rate proportional to $\langle m_\nu \rangle^2_{LL}$, in this limit there is a close analogy between the matrix elements governing 0ν and 2ν decay.

Before making a connection between experiment and the parameters governing lepton number violation, these nuclear matrix elements must be evaluated. The approach followed by our group in Los Alamos[5] is based on the nuclear shell model. My discussion here of the nuclear physics will be only qualitative.

The similarity between some 0ν matrix elements and those mediating 2ν decay suggests a natural check on calculated 0ν decay rates: do the nuclear wave functions reproduce known 2ν decay rates. The results shown in the table are surprising. Theory predicts a suppressed rate for ^{48}Ca, in agreement with the measurement of the Columbia group[6], and a rate for ^{82}Se near that found in the cloud chamber experiment of Moe and Lowenthal.[7] However, the three geochemical total $\beta\beta$ decay rates are significantly slower than the theoretical 2ν rate estimates. The most alarming discrepancy exists for the Te isotopes, with the theoretical and experimental absolute rates differing by two order of magnitude. In the case of ^{82}Se two experimental techniques (cloud chamber measurements and geochemical determinations) give results that disagree by as much as a factor of 25. The theoretical rate, though closer to the cloud chamber result, is bracketed by these measurements.

Table 1: Calculated and Experimental Double Gamow-Teller Matrix Elements M_{GT}

Nucleus	$\lvert M_{GT}\rvert_{theory}$ [5]	$\lvert M_{GT}\rvert_{exp}$
^{48}Ca	0.22	≤ 0.20 [6]
^{76}Ge	1.28	
^{82}Se	0.94	1.51 [7]
		0.40 [8]
		0.29 [9]
^{128}Te	1.47	0.21-0.25 [10]
		<0.19 [11]
^{130}Te	1.48	0.11-0.14 [11]
		0.19 [10]

Why have such large discrepancies between theory and geochemistry received so little attention previously? One point overlooked until recently is the enhancement of the phase space by relativistic effects in the Coulomb distortion of the outgoing electron waves[5]. This increases the theoretical rates for the Te isotopes by factors of 4 to 5. A second and perhaps more important point is the tendency for naive nuclear calculations to underestimate decay rates. Consider the double Gamow-Teller operator (Eq. (4a)). In a simple description for the Te isotopes one might restrict the valence neutron holes to the $1h_{11/2}$ shell and the valence protons to the $1g_{7/2}$ and $2d_{5/2}$ shells. Because the valence protons and neutron holes have different orbital angular momenta, $M_{GT}=0$. Clearly one should require of a realistic calculation that the shell model basis for the final state include those configurations that can connect to the initial state via the double Gamow-Teller operator, and conversely. If this condition is satisfied the prediction for M_{GT} should depend on the character of the nuclear force, not on spurious effects associated with basis truncation.

Calculations attempting to respect this "sum rule" constraint have been performed with the shell model,[5] with the random phase approximation,[12] and with the Nilsson-pairing model.[13] The results are quite similar. The Nilsson-pairing calculations demonstrate explicitly that the nuclear pairing force is responsible for the strong matrix elements derived in the shell model. Although one can obtain matrix elments in agreement with geochemistry by significantly weakening this component of the nuclear force, such an adjustment runs contrary to prevailing prejudice. Note that this dependence of $\beta\beta$ decay on the pairing force contrasts with that of single β decay, where transition rates tend to diminish as the strength of the pairing force is increased.

Finally, a remaining worry is the use of the closure approximation in place of the inverse-energy-weighted sum over virtual intermediate nuclear states. It is difficult to avoid this approximation in calculations of sufficient complexity to be otherwise realistic.

The uncertainties stemming from the three-sided discrepancy between nuclear theory, geochemistry, and laboratory experiments clearly affect the confidence we have in $\beta\beta$ decay limits on lepton number violation. One way of parameterizing this uncertainty depends on a relation suggested by Primakoff and Rosen[14])

$$\frac{\langle 0_f^+ | \frac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\vec{\sigma}(i)\cdot\vec{\sigma}(j) \frac{g(r_{ij})}{r_{ij}} | 0_i^+ \rangle}{\langle 0_f^+ | \frac{1}{2} \sum_{ij} \tau_+(i)\tau_+(j)\vec{\sigma}(i)\cdot\vec{\sigma}(j) | 0^+ \rangle} \sim \frac{1}{R_o}$$

where R_o is the nuclear radius. For the shell model calculations in ^{76}Ge, ^{82}Se, ^{128}Te, and ^{130}Te this proportionality between matrix elements holds rather well, though with a somewhat different strength ($R_o \rightarrow (0.58 \pm 0.03) R_o$). Thus a reasonable procedure may be to scale all theoretical matrix elements by a common factor chosen to fit a given experimental 2ν rate. The scale factors determined from conflicting experiments provide a measure of the associated uncertainty in the limits on lepton number violation.

Those limits imposed by the ^{76}Ge and ^{82}Se 0ν $\beta\beta$ decay results, $\tau_{1/2}^{0\nu} \lesssim 5\cdot 10^{22}$ y[15]) and $\lesssim 3.1\cdot 10^{21}$ y,[16]) respectively, are shown in Figs. 3. and 4. The dashed lines in Fig. 4 correspond to scaled calculations.

Fig. 3 employs theoretical matrix elements only because the ^{76}Ge 2ν decay rate is not known. (Note that Prof. Fiorini is reporting an improved bound on $\tau^{0\nu}_{1/2}$ for ^{76}Ge at this meeting.)

Figure 3

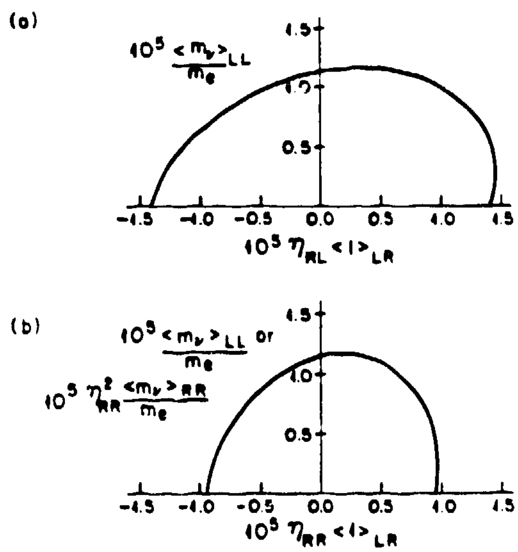
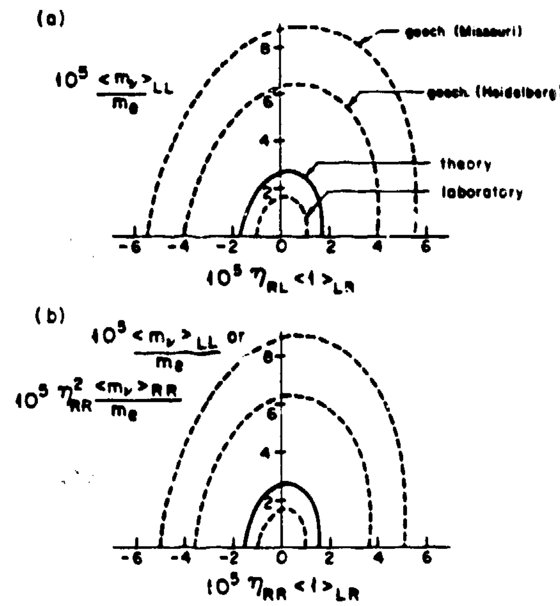


Figure 4



It may be possible to circumvent the question of matrix element normalization in one case. Pontecorvo,[17] in his discussion of $\Delta\ell=2$ superweak interactions, suggested that, as the $\beta\beta$ decay matrix elements for the neighboring nuclei ^{128}Te and ^{130}Te are likely similar, the ratio of these decay rates may depend primarily on phase space. With this assumption one finds

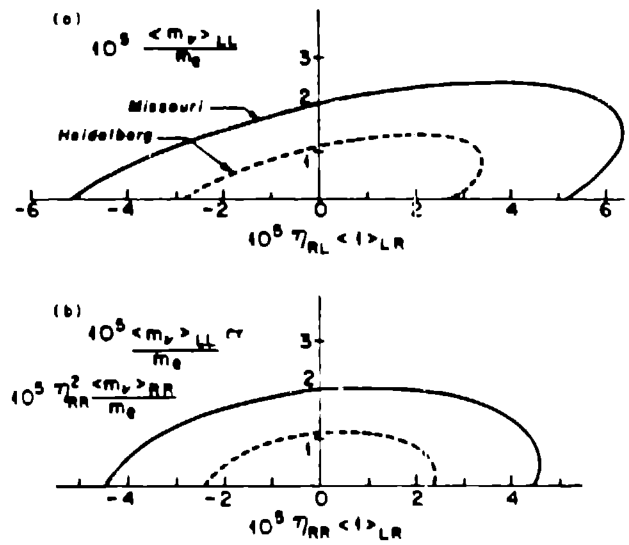
$$\frac{\tau_{1/2}(128)}{\tau_{1/2}(130)} = \begin{cases} 5130, & 2\nu \text{ decay} \\ 25, & 0\nu \text{ decay with } \langle m_\nu \rangle_{LL} \neq 0 \\ 116, & 0\nu \text{ decay with } \eta_{RL} \langle 1 \rangle_{LR} \neq 0 \end{cases}$$

The experimental result of the Missouri group[10)]

$$\frac{\tau_{1/2}(128)}{\tau_{1/2}(130)} = 1580$$

is bracketed by the 2ν and 0ν ratios and, as Bryman and Picciotto suggested,[18)] may indicate that both modes contribute. Some choices for the parameters governing lepton number violation that would generate this ratio are given by the solid curves in Fig. 5.

Figure 5



A recent Heidelberg result[11], however, conflicts with the Missouri value

$$\frac{\tau_{1/2}(128)}{\tau_{1/2}(130)} > 3040, \ 95\% \text{ C.L.}$$

This value is consistent with 2ν decay and imposes stringent bounds on lepton number violation, as shown by the dashed curves in Fig. 5.

In summary, the results in Figs. 3, 4, and 5 show $\langle m_\nu \rangle_{LL} \lesssim (4-19)$ eV, with the range reflecting the experimental and theoretical uncertainties discussed previously. One also finds $|\eta_{RL} \langle 1 \rangle_{LR}| \lesssim (1-5) \cdot 10^{-5}$. A recent interpretation of the ITEP tritium β decay measurements[19] suggested $m_\nu \sim 31 \pm 2$ eV and, as a bound independent of atomic physics, $m_\nu \gtrsim 20$ eV. Thus, pending verification of these results, present $\beta\beta$ decay studies could rule out the possibility that the electron neutrino is a Majorana mass eigenstate. The ITEP results have also raised speculations about multiple mass eigenstates. One possibility, $m_1 \sim 22$ eV (95%) and $m_2 \sim 114$ eV (5%), would yield a $\beta\beta$ decay mass of $\langle m_\nu \rangle_{LL} \sim$

15 eV for Majorana mass eigenstates with opposite CP. This result may be marginally consistent with limits on 0ν $\beta\beta$ decay and, because of the small mixing angle, with bounds on neutrino oscillations (see Prof. Boehm's talk at this meeting).

If heavy Majorana neutrinos couple to the electron, the appropriate propagator is

$$\frac{g(r_{ij})}{r_{ij}} \rightarrow \frac{e^{-m_\nu r_{ij}}}{r_{ij}} \tag{12}$$

If both light and heavy neutrinos contribute to the 0ν decay amplitude, the different radial dependence of the propagators prevents a simple factorization of the nuclear physics from the particle physics couplings. Some perverse consequences of this have been discussed by Halprin, Petcov, and Rosen[20] and by our group at Los Alamos.[1].

If only heavy neutrinos with masses large compared to the inverse size of the nucleon contribute, matters simplify. One finds[1),21)]

$$\left\langle\frac{e^{-m_\nu r_{ij}}}{r_{ij}}\right\rangle \sim \left(\frac{M_A}{m_\nu}\right)^2 \left\langle\frac{1}{r_{ij}} F(M_A r_{ij})\right\rangle \tag{13}$$

where M_A is the nucleon form factor mass. Thus, again only a common radial integral arises. The nucleon finite size leads to a soft $1/m_\nu^2$ dependence of the $\beta\beta$ decay matrix elements. A rough estimate of the limits imposed by the ^{76}Ge 0ν $\beta\beta$ decay half life on the coupling of a single massive Majorana neutrino is given in Fig. 6. For $m_\nu \sim 1$ GeV one finds $|U_e^L|^2 \lesssim 10^{-6}$.

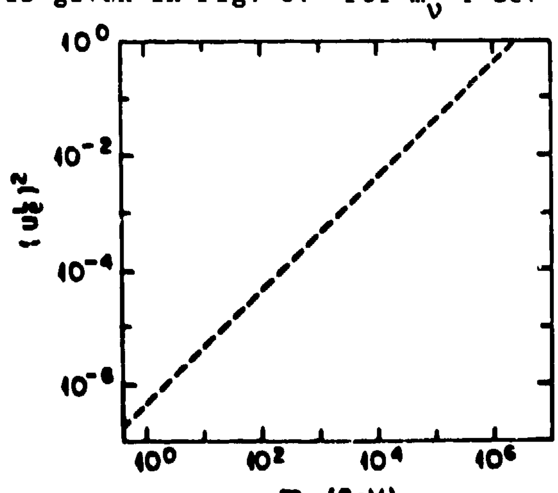


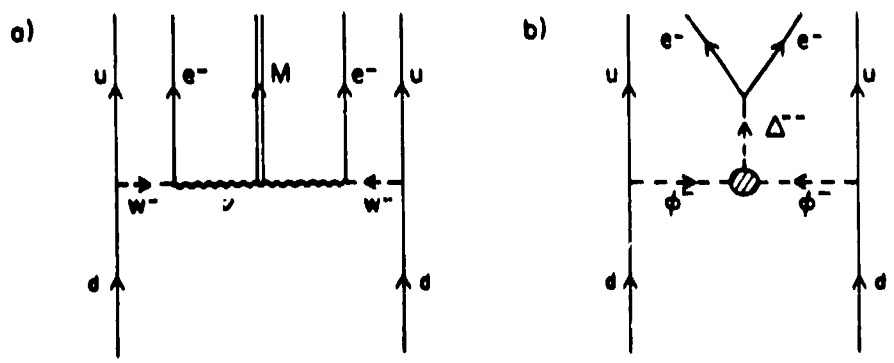
Figure 6: Limits on the coupling of a heavy neutrino

Interest in $0^+ \rightarrow 2^+$ 0ν $\beta\beta$ decay was stimulated by the discussions of Rosen [22] and Doi et al.[23]. Ths process can only occur via the right-handed couplings η_{RL} and η_{RR}. Calculations[1] indicate that the constraints on these couplings imposed by bounds on $0^+ \rightarrow 2^+$ decay are much less stringet than those from $0^+ \rightarrow 0^+$ decay. Of course, if 0ν $\beta\beta$ decay is discovered, this process could prove of great value in distinguishing neutrino mass effects from those of right-handed currents. An additional reason for interest in $0^+ \rightarrow 2^+$ decays is the existence of single hadron mechanisms (e.g., $n \leftrightarrow \Delta^{++}$).

Various authors have recently discussed several exotic mechanisms for 0ν $\beta\beta$ decay. Estimates of the importance of the Higgs exchange mechanism (Fig. 7b), first considered by Mohapatra and Vergados[24], have varied widely depending on specific assumptions in the nuclear and particle physics. Neutrinoless $\beta\beta$ decay accompanied by Majoran production has been discussed by Georgi, Glashow, and Nussinov[25] (Fig. 7a). This process would thwart conventional searches for 0ν decay because it does not produce the characteristic spectrum of electron pairs with a fixed total energy. One interesting consequence of this mechanism is the occurrence of neutrinoless double electron capture in the absence of an accidental energy balance between the initial and final states.

I thank S. P. Rosen, G. J. Stephenson, Jr., and D. Strottman for many interesting discussions and collaborations.

Figure 7: Majorana and Higgs exchange mechanisms.



References

1) For general reviews see W. C. Haxton and G. J. Stephenson, Jr., Los Alamos preprint LA-UR-84-396 (to be published in Progress in Particle and Nuclear Physics); S. P. Rosen, in Science Underground, AIP Conf. Proc. 96, ed. M. M. Nieto et al. (New York, 1983); M. Doi, T. Kotani, H. Nishiura, and E. Takasugi, Prog. Theor. Phys. 69 (1983) 602.

2) G. Racah, Nuovo Cimento 14 (1937) 327.

3) C. S. Wu, E. Ambler, R. W. Hayward, D. D. Hopper, and R. D. Hudson, Phys. Rev. 105 (1957) 1413; R. Garwin, L. Lederman, and M. Weinrich, Phys. Rev. 105 (1957) 1415; J. I. Freidman and V. L. Telegdi, Phys. Rev. 105 (1957) 1681; T. D. Lee and C. N. Yang, Phys. Rev. 104 (1956) 254.

4) M. Gell-Mann, P. Ramond, and R. Slansky, unpublished and in "Supergravity", ed. by P. van Nieuwenhuizen and D. Z. Freedman (North Holland, Amsterdam, 1979).

5) W. C. Haxton, G. J. Stephenson, Jr., and D. Strottman, Phys. Rev. Lett. 47 (1981) 153 and Phys. Rev. D25 (1982) 2360; W. C. Haxton, S. P. Rosen, and G. J. Stephenson, Jr., Phys. Rev. D26 (1982) 1805.

6) R. K. Bardin, P. J. Gollon, J. D. Ullman, and C. S. Wu, Phys. Lett. 26B (1967) 112 and Nucl. Phys. A158 (1970) 337.

7) M. K. Moe and D. D. Lowenthal, Phys. Rev. C22 (1980) 2186.

8) T. Kirsten, in Science Underground, AIP Conf. Proc. 96, ed. M. M. Nieto et al. (New York, 1983).

9) B. Srinivasan, E. C. Alexander, Jr., R. D. Beaty, D. E. Sinclair, and O. K. Manuel, Econ. Geo. 68 (1973) 252.

10) E. W. Hennecke, O. K. Manuel, and D. D. Sabu, Phys. Rev. C11 (1975) 1378; E. W. Hennecke, Phys. Rev. C17 (1978) 1168.

11) T. Kirsten, H. Richter, and E. Jessberger, Phys. Rev. Lett. 50 (1983) 474 and Z. Physik C16 (1983) 189 and references therein.

12) A. H. Huffman, Phys. Rev. C2 (1970) 742.

13) L. Zamick and N. Auerbach, Phys. Rev. C26 (1982) 2185.

14) H. Primakoff and S. P. Rosen, Rep. Prog. Phys. 22 (1959) 121 and Proc. Phys. Soc. 78 (1961) 464.

15) E. Bellotti, E. Fiorini, C. Liguori, A. Pallia, A Sarracino and L. Zanotti, Lett. Nuovo Cimento 33 (1982) 273. Also see F. T. Avignone III, R. L. Brodzinski, D. P. Brown, J. C. Evans, Jr., W. K. Hensley, J. H. Reeves, and W. A. Wogman, Phys. Rev. Lett. 50 (1983) 721 and A. Forster, H. Kwon, J. K. Markey, F. Boehm, and H. E. Henrikson, talk presented at Conf. on Low Energy Tests of Conservation Laws in Particle Physics, Blacksburg, 1983.

16) B. T. Cleveland et al., Phys. Rev. Lett. 35 (1975) 757.

17) B. Pontecorvo, Phys. Lett 26B (1968) 630.

18) D. Bryman and C. Picciotto, Rev. Mod. Phys. 50 (1978) 11.

19). S. Boris, A. Golutvin, L. Laptin, V. Lubimov, V. Nagovizin, E. Novikov, V. Nozik, V. Soloshenko, I. Tichomirov, and E. Tretjakov, talk presented at the Int. Europhysics Conf. on High Energy Physics, Brighton, 1983.

20) A. Halprin, S. T. Petcov, and S. P. Rosen, Phys. Lett. 125B (1983) 335.

21) J. D. Vergados, Univ. of Washington preprint (1984).

22) S. P. Rosen, talk presented at Orbis Scientiae, Univ. of Miami (1981).

23) M. Doi, T. Kotani, H. Nishiura, K. Okuda, and E. Takasugi, Phys. Lett. 103B (1981) 219 and 113B (1982) E513.

24) R. N. Mohapatra and J. D. Vergados, Phys. Rv. Lett. 47 (1981) 1713.

25) H. M. Georgi, S. L. Glashow, and S. Nussinov, Nucl. Phys. B193 (1981) 297.